

IASSNS-HEP-94/64

February 1, 2008

T

## Custodial nonabelian gauge symmetries in realistic superstring derived models

ALON E. FARAGGI<sup>†</sup>

School of Natural Sciences

Institute for Advanced Study

Olden Lane

Princeton, NJ 08540

### ABSTRACT

A well known problem in supersymmetric models is the presence of, lepton and baryon number violating, dimension four operators. The traditional R parity solution may not be suitable if one tries to incorporate supersymmetry into a Planck scale theory. I propose a different solution in the context of realistic string models. I show that realistic string models can give rise to custodial non-abelian gauge symmetries under which only the leptons or quarks transform. I explain how such symmetries arise in a class of free fermionic models that are based on  $Z_2 \times Z_2$  orbifold with standard embedding. The custodial symmetries forbid proton decay from dimension four operators while allowing R parity violation.

---

<sup>†</sup> Work supported by an SSC fellowship. e-mail address: faraggi@sns.ias.edu

Supersymmetry [1] is a phenomenologically appealing extension of the Standard Model. While LEP precision data strongly constrain other attractive extensions of the Standard Model, supersymmetry below the TeV scale is in good agreement with all experiments to date. Moreover, it is derived from superstring theory and provides a solution to the gauge hierarchy problem. However, despite this attractive properties, supersymmetry gives rise to dimension four, baryon and lepton violating, operators that result in fast proton decay. The dangerous dimension four operators are forbidden by postulating an extra matter parity symmetry. In superstring theory [2] the problem is more severe because such a symmetry cannot be imposed at will. Early in the days of string inspired phenomenology it was pointed out that matter parity was not automatic in most string-derived models and that renormalizable dimension four operators were present in generic string vacua [3]. By going to a particular symmetric vacuum, these operators could often be avoided [4]. However to produce a realistic low energy mass spectrum it is, in general, necessary to perturb away from the symmetric points in moduli space. When perturbing away from the symmetric point, it is difficult to envision how one can control the absence of dimension four operators.

In  $SO(10)$  based models absence of a cubic sixteen operator forbid the dimension four operators. However, the dangerous dimension four operators may still be induced from quartic sixteen operators, if one of the spinorial sixteen of  $SO(10)$  gets a GUT or Planck scale VEV, as such a VEV breaks matter parity. In terms of standard model multiplets the dangerous operators are

$$\eta_1(u_L^C d_L^C d_L^C N_L^c)\Phi + \eta_2(d_L^C Q L N_L^c)\Phi, \quad (1)$$

where  $N_L^c$  is the Standard Model singlet in the 16 of  $SO(10)$ .  $\Phi$  is a combination of fields which fixes the string selection rules and gets a VEV of  $O(M/10)$ , where  $M = M_{Pl}/2\sqrt{8\pi}$ . From Eq. (1), it is seen that the ratio  $\langle N_L^c \rangle/M$  controls the rate of proton decay. A search through nonrenormalizable terms shows that terms of the form of Eq. (1) are, in general, generated in string models [5,6]. An additional

$U(1)_{Z'}$  that is unbroken down to low energies suppresses the dangerous dimension four operators in  $SO(10)$  based models. The problem with this solution is that often the scale of  $U(1)_{Z'}$  breaking is the scale of a see-saw mechanism that suppresses the left-handed neutrino masses. While in field theory this scale can be near the TeV scale, in superstring models, because the terms in the seesaw mass matrix are usually obtained from nonrenormalizable terms, the scale of  $U(1)_{Z'}$  breaking is often required to be much higher [9].

In this paper I propose that in string theory a different mechanism is required. In this mechanism custodial nonabelian symmetries are obtained under which only the leptons or quarks transform. The custodial symmetry then allows only baryon or lepton violating dimension four operators but not both. Therefore with the custodial nonabelian symmetries R-parity may be broken close to the Planck scale but proton decay from dimension four operators is suppressed. I construct a toy model to illustrate this mechanism. In this model the Standard Model leptons transform under a custodial  $SU(2)$  symmetry. In the toy model the custodial symmetries arise due to additional space-time vector bosons that are obtained from twisted sectors. These twisted sectors are obtained from boundary condition vectors that break the  $SO(10)$  symmetry and correspond to “Wilson lines” in the orbifold formulation. Contrary to the gauged  $B - L$  symmetry that arises in string models solely from the world-sheet gauge degrees of freedom [7,8], the custodial symmetries arise from mixture of the gauge and internal degrees of freedom. I discuss some additional aspects of similar extended symmetries and their possible phenomenological implications.

The superstring models that I present are constructed in the free fermionic formulation [10]. In this formulation all the degrees of freedom needed to cancel the conformal anomaly are represented in terms of internal free fermions propagating on the string world-sheet. Under parallel transport around a noncontractible loop, the fermionic states pick up a phase. Specification of the phases for all world-sheet fermions around all noncontractible loops contributes to the spin structure

of the model. The possible spin structures are constrained by string consistency requirements (e.g. modular invariance). A model is constructed by choosing a set of boundary condition vectors, which satisfies the modular invariance constraints. The basis vectors,  $b_k$ , span a finite additive group  $\Xi = \sum_k n_k b_k$  where  $n_k = 0, \dots, N_{z_k} - 1$ . The physical massless states in the Hilbert space of a given sector  $\alpha \in \Xi$ , are obtained by acting on the vacuum with bosonic and fermionic operators and by applying the generalized GSO projections. The  $U(1)$  charges,  $Q(f)$ , with respect to the unbroken Cartan generators of the four dimensional gauge group, which are in one to one correspondence with the  $U(1)$  currents  $f^* f$  for each complex fermion  $f$ , are given by:

$$Q(f) = \frac{1}{2}\alpha(f) + F(f) \quad (2)$$

where  $\alpha(f)$  is the boundary condition of the world-sheet fermion  $f$  in the sector  $\alpha$ , and  $F_\alpha(f)$  is a fermion number operator counting each mode of  $f$  once (and if  $f$  is complex,  $f^*$  minus once). For periodic fermions,  $\alpha(f) = 1$ , the vacuum is a spinor in order to represent the Clifford algebra of the corresponding zero modes. For each periodic complex fermion  $f$  there are two degenerate vacua  $|+\rangle, |-\rangle$ , annihilated by the zero modes  $f_0$  and  $f_0^*$  and with fermion numbers  $F(f) = 0, -1$ , respectively.

The realistic models in the free fermionic formulation are generated by a basis of boundary condition vectors for all world-sheet fermions [8,11–15]. The basis is constructed in two stages. The first stage consist of the NAHE set [12], which is a set of five boundary condition basis vectors,  $\{\mathbf{1}, S, b_1, b_2, b_3\}$ . The gauge group after the NAHE set is  $SO(10) \times SO(6)^3 \times E_8$  with  $N = 1$  space-time supersymmetry. The vector  $S$  is the supersymmetry generator and the superpartners of the states from a given sector  $\alpha$  are obtained from the sector  $S + \alpha$ . The space-time vector bosons that generate the gauge group arise from the Neveu-Schwarz sector and from the sector  $1 + b_1 + b_2 + b_3$ . The Neveu-Schwarz sector produces the generators of  $SO(10) \times SO(6)^3 \times SO(16)$ . The sector  $1 + b_1 + b_2 + b_3$  produces the spinorial 128 of  $SO(16)$  and completes the hidden gauge group to  $E_8$ . The vectors  $b_1, b_2$  and  $b_3$

correspond to the three twisted sectors in the corresponding orbifold formulation and produce 48 spinorial 16 of  $SO(10)$ , sixteen from each sector  $b_1$ ,  $b_2$  and  $b_3$ .

The NAHE set divides the 44 right-moving and 20 left-moving real internal fermions in the following way:  $\bar{\psi}^{1,\dots,5}$  are complex and produce the observable  $SO(10)$  symmetry;  $\bar{\phi}^{1,\dots,8}$  are complex and produce the hidden  $E_8$  gauge group;  $\{\bar{\eta}^1, \bar{y}^{3,\dots,6}\}$ ,  $\{\bar{\eta}^2, \bar{y}^{1,2}, \bar{\omega}^{5,6}\}$ ,  $\{\bar{\eta}^3, \bar{\omega}^{1,\dots,4}\}$  give rise to the three horizontal  $SO(6)$  symmetries. The left-moving  $\{y, \omega\}$  states are divided to,  $\{y^{3,\dots,6}\}$ ,  $\{y^{1,2}, \omega^{5,6}\}$ ,  $\{\omega^{1,\dots,4}\}$ . The left-moving  $\chi^{12}, \chi^{34}, \chi^{56}$  states carry the supersymmetry charges. Each sector  $b_1$ ,  $b_2$  and  $b_3$  carries periodic boundary conditions under  $(\psi^\mu | \bar{\psi}^{1,\dots,5})$  and one of the three groups:  $(\chi_{12}, \{y^{3,\dots,6} | \bar{y}^{3,\dots,6}\}, \bar{\eta}^1)$ ,  $(\chi_{34}, \{y^{1,2}, \omega^{5,6} | \bar{y}^{1,2} \bar{\omega}^{5,6}\}, \bar{\eta}^2)$  and  $(\chi_{56}, \{\omega^{1,\dots,4} | \bar{\omega}^{1,\dots,4}\}, \bar{\eta}^3)$ . The division of the internal fermions is a reflection of the underlying  $Z_2 \times Z_2$  orbifold compactification [16]. The set of internal fermions  $\{y, \omega | \bar{y}, \bar{\omega}\}^{1,\dots,6}$  corresponds to the left-right symmetric conformal field theory of the heterotic string, or to the six dimensional compactified manifold in a bosonic formulation. This set of left-right symmetric internal fermions plays a fundamental role in the determination of the low energy properties of the realistic free fermionic models.

The second stage of the basis construction consist of adding three additional basis vectors to the NAHE set. The three additional basis vectors correspond to “Wilson lines” in the orbifold formulation. Three additional vectors are needed to reduce the number of generations to three, one from each sector  $b_1$ ,  $b_2$  and  $b_3$ . The additional basis vectors distinguish between different models and determine their low energy properties. The allowed boundary conditions in the additional basis vectors are constrained by the string consistency constraints, i.e. modular invariance and world-sheet supersymmetry. The choice of boundary conditions to the set of internal fermions  $\{y, \omega | \bar{y}, \bar{\omega}\}^{1,\dots,6}$  determines the low energy properties, like the number of generations, Higgs doublet-triplet splitting and Yukawa couplings. The low energy phenomenological requirements impose strong constraints on the possible assignment of boundary conditions to the set of of internal world-sheet fermions  $\{y, \omega | \bar{y}, \bar{\omega}\}^{1,\dots,6}$  [12].

For some choices of the additional basis vectors that extend the NAHE set, there exist a combination

$$X = n_\alpha \alpha + n_\beta \beta + n_\gamma \gamma \quad (3)$$

for which  $X_L \cdot X_L = 0$  and  $X_R \cdot X_R \neq 0$ . Such a combination may produce additional space-time vector bosons, depending on the choice of GSO phases. For example, in the model of Ref. [13] the combination  $X = b_1 + b_2 + b_3 + \alpha + \beta + \gamma$  has  $X_L \cdot X_L = 0$  and  $X_R \cdot X_R = 8$ . The space-time vector bosons from this sector are projected out by the choice of GSO phases, and this vector combination produces only space-time scalar supermultiplets. On the other hand in the model of Ref. [14] with the modified GSO phase

$$c \begin{pmatrix} \gamma \\ 1 \end{pmatrix} = +1 \rightarrow c \begin{pmatrix} \gamma \\ 1 \end{pmatrix} = -1$$

additional space-time vector bosons are obtained from the sector  $1 + \alpha + 2\gamma$ . The gauge group after applying the GSO projections is  $SU(3)_C \times SU(2)_L \times SU(2)_c \times U(1)_{C'} \times U(1)_L \times U(1)^5 \times SU(5) \times SU(3) \times U(1)$ .

The gauge group arises as follows: the NS sector produces the generators of  $SU(3)_C \times SU(2)_L \times U(1)_C \times U(1)_L \times U(1)_{1,2,3} \times U(1)_{4,5,6} \times SU(3) \times SO(4) \times U(1)_H \times U(1)_{7,8,9}$  where

$$U(1)_C = Tr U(3)_C \Rightarrow Q_C = \sum_{i=1}^3 Q(\bar{\psi}^i), \quad (4a)$$

$$U(1)_L = Tr U(2)_L \Rightarrow Q_L = \sum_{i=4}^5 Q(\bar{\psi}^i), \quad (4b)$$

$$U(1)_H = Tr U(3)_H \Rightarrow Q_H = \sum_{i=5}^7 Q(\bar{\phi}^i). \quad (4c)$$

$U(1)_{1,2,3}$ ,  $U(1)_{4,5,6}$  and  $U(1)_{5,6,7}$  arise from the world-sheet currents  $\bar{\eta}^i \bar{\eta}^{i*}$  ( $i = 1, 2, 3$ ),  $\bar{y}^3 \bar{y}^6$ ,  $\bar{y}^1 \bar{\omega}^5$ ,  $\bar{\omega}^2 \bar{\omega}^4$ , and  $\bar{\phi}^1 \bar{\phi}^{1*}$ ,  $\bar{\phi}^2 \bar{\phi}^{2*}$ ,  $\bar{\phi}^8 \bar{\phi}^{8*}$ , respectively. The sector  $1 +$

$b_1 + b_2 + b_3$  produces the representations  $(3, 2)_{-5} \oplus (\bar{3}, 2)_5$  and  $2_{-3} \oplus 2_3$  of  $SU(3) \times SU(2)_r \times U(1)_{h5}$  and  $SU(2)_\ell \times U(1)_{h3}$  respectively, where  $SU(2)_r \times SU(2)_\ell$  are the two  $SU(2)$ 's in the isomorphism  $SO(4) \sim SU(2)_r \times SU(2)_\ell$ . Thus, the  $E_8$  symmetry reduces to  $SU(5) \times SU(3) \times U(1)^2$ . The  $U(1)$ 's in  $SU(5)$  and  $SU(3)$  are given by  $U(1)_{h5} = -3U_7 + 3U_8 + U_H - 3U_9$  and  $U(1)_{h3} = U_7 + U_8 + U_H + U_9$  respectively. The remaining  $U(1)$  symmetries in the hidden sector,  $U(1)_{7'}$  and  $U(1)_{8'}$ , correspond to the world-sheet currents  $\bar{\phi}^1 \bar{\phi}^{1*} - \bar{\phi}^8 \bar{\phi}^{8*}$  and  $-2\bar{\phi}^j \bar{\phi}^{j*} + \bar{\phi}^1 \bar{\phi}^{1*} + 4\bar{\phi}^2 \bar{\phi}^{2*} + \bar{\phi}^8 \bar{\phi}^{8*}$  respectively, where summation on  $j = 5, \dots, 7$  is implied.

The sector  $1 + \alpha + 2\gamma$  produces two additional space-time vector bosons, which are singlets of the nonabelian group but carry  $U(1)$  charges. One combination of the  $U(1)$  symmetries

$$U_C + U_4 + U_5 + U_6 + U_{7'} \quad (5)$$

is the  $U(1)$  of the custodial  $SU(2)$  symmetry. The two space-time vector bosons from the sector  $1 + \alpha + 2\gamma$  produce the two additional vector bosons of the custodial  $SU(2)$  gauge group. The remaining orthogonal combinations are

$$U_{C'} = \frac{1}{3}U_C - \frac{1}{2}U_{7'}, \quad (6a)$$

$$U_{4'} = U_4 - U_6, \quad (6b)$$

$$U_{5'} = U_4 + U_5 - 2U_6, \quad (6c)$$

$$U_{7''} = U_C - \frac{5}{3}(U_4 + U_5 + U_6) + U_{7'}. \quad (6d)$$

The full massless spectrum now transforms under the final gauge group,  $SU(3)_C \times SU(2)_L \times SU(2)_c \times U(1)_{C'} \times U(1)_L \times U(1)_{1,2,3} \times U(1)_{4'} \times U(1)_{5'} \times U(1)_{7''} \times U(1)_8$ .

(a) The Neveu-Schwarz  $O$  sector gives, in addition to the graviton, dilaton, antisymmetric tensor and spin 1 gauge bosons, the following scalar representations:

$$h_1 \equiv [(1, 0); (2, -1)]_{1,0,0,0,0,0} \quad \Phi_{23} \equiv [(1, 0); (1, 0)]_{0,1,-1,0,0,0} \quad (7a, b)$$

$$h_2 \equiv [(1, 0); (2, -1)]_{0,1,0,0,0,0} \quad \Phi_{13} \equiv [(1, 0); (1, 0)]_{1,0,-1,0,0,0} \quad (7c, d)$$

$$h_3 \equiv [(1, 0); (2, -1)]_{0,0,1,0,0,0} \quad \Phi_{12} \equiv [(1, 0); (1, 0)]_{1,-1,0,0,0,0} \quad (7e, f)$$

(and their conjugates  $\bar{h}_1$  etc.). Finally, the Neveu–Schwarz sector gives rise to three singlet states that are neutral under all the  $U(1)$  symmetries.  $\xi_{1,2,3} : \chi_{\frac{1}{2}}^{12} \bar{\omega}_{\frac{1}{2}}^3 \bar{\omega}_{\frac{1}{2}}^6 |0\rangle_0$ ,  $\chi_{\frac{1}{2}}^{34} \bar{y}_{\frac{1}{2}}^5 \bar{\omega}_{\frac{1}{2}}^1 |0\rangle_0$ ,  $\chi_{\frac{1}{2}}^{56} \bar{y}_{\frac{1}{2}}^2 \bar{y}_{\frac{1}{2}}^4 |0\rangle_0$ .

(b) The  $S + b_1 + b_2 + \alpha + \beta$  sector gives

$$h_{45} \equiv [(1, 0); (2, -1)]_{\frac{1}{2}, \frac{1}{2}, 0, 0, 0, 0} \quad h'_{45} \equiv [(1, 0); (2, -1)]_{-\frac{1}{2}, -\frac{1}{2}, 0, 0, 0, 0} \quad (8a, b)$$

$$\Phi_{45} \equiv [(1, 0); (1, 0)]_{-\frac{1}{2}, -\frac{1}{2}, -1, 0, 0, 0} \quad \Phi'_{45} \equiv [(1, 0); (1, 0)]_{-\frac{1}{2}, -\frac{1}{2}, 1, 0, 0, 0} \quad (8c, d)$$

$$\Phi_1 \equiv [(1, 0); (1, 0)]_{-\frac{1}{2}, \frac{1}{2}, 0, 0, 0, 0} \quad \Phi_2 \equiv [(1, 0); (1, 0)]_{-\frac{1}{2}, \frac{1}{2}, 0, 0, 0, 0} \quad (8e, f)$$

(and their conjugates  $\bar{h}_{45}$ , etc.). The states are obtained by acting on the vacuum with the fermionic oscillators  $\bar{\psi}^{4,5}, \bar{\eta}^3, \bar{y}_5, \bar{\omega}_6$ , respectively (and their complex conjugates for  $\bar{h}_{45}$ , etc.).

The sectors  $b_j \oplus 1 + \alpha + 2\gamma$  produce the three light generations, one for each of the sectors  $b_j$  ( $j = 1, 2, 3$ ). The states from these sectors and their decomposition under the entire gauge group are shown in table 1. From table 1 we see that only the lepton supermultiplets,  $\{L, e_L^c, N_L^c\}$  transform as doublets under the custodial  $SU(2)$  gauge group while the quarks are singlets. The remaining matter states in the massless spectrum and their quantum numbers are given in table 2.

The model contains three anomalous  $U(1)$  symmetries:  $\text{Tr}U_1 = 24$ ,  $\text{Tr}U_2 = 24$ ,  $\text{Tr}U_3 = 24$ . Of the three anomalous  $U(1)$ s, two can be rotated by an orthogonal transformation. One combination remains anomalous and is uniquely given by:  $U_A = k \sum_j [\text{Tr}U(1)_j] U(1)_j$ , where  $j$  runs over all the anomalous  $U(1)$ s. For convenience, I take  $k = 1/24$ . Therefore, the anomalous combination is given by:

$$U_A = U_1 + U_2 + U_3, \quad \text{Tr}Q_A = 72. \quad (9a)$$

The two orthogonal combinations are not unique. Different choices are related by orthogonal transformations. One choice is given by:

$$U'_1 = U_1 - U_2, \quad U'_2 = U_1 + U_2 - 2U_3. \quad (9b, c)$$



Together with the other anomaly free  $U(1)$ s, they are free from gauge and gravitational anomalies. The cancelation of all mixed anomalies among the  $U(1)$ s is a non trivial consistency check of the massless spectrum of the model. The “anomalous”  $U(1)_A$  is broken by the Dine-Seiberg-Witten mechanism [18] in which some states in the massless spectrum obtain nonvanishing VEVs that cancel the anomalous  $U(1)$  D-term equation. A particular example, in the model under consideration, is given by the set  $\{\Phi_{45}, \Phi'_{45}\}$  with  $|\langle\Phi_{45}\rangle|^2 = 3|\langle\Phi'_{45}\rangle|^2 = \frac{3g^2}{16\pi^2}$ .

Nonvanishing VEVs of the states form the sectors  $b_j + 2\gamma$  break the  $U(1)$  symmetries to  $U(1)_C \times U(1)_L$ . The VEV of  $N_L^c$  breaks the custodial  $SU(2)_c$  symmetry and the remaining  $U(1)$  symmetry. The surviving combination  $1/3U(1)_C + 1/2U(1)_L$  is the Standard Model weak hypercharge. Only the leptons  $L_j$ ,  $e_j$  and  $N_j$  transform as doublets under the custodial  $SU(2)_c$  gauge group whereas the quarks are singlets of  $SU(2)_c$  (see table 1). Consequently, the term  $QLd_L^c N_L^c \Phi$  is allowed while the term

$$u_L^C d_L^C d_L^C N_L^c \Phi \quad (10)$$

is forbidden due to invariance under  $SU(2)_c$ . Thus, the VEV of  $N_L^c$  can be of order  $M_{Pl}$ , and although it breaks matter parity, it does not imply proton decay from dimension four operators. While the lepton number violating dimension four operator,  $QLd^c$ , is allowed and may be unsuppressed, baryon number violating dimension four operators are forbidden. An important implication of R parity violation is that the lightest supersymmetric particle is unstable. Analysis of models that allow this type of matter parity breaking has been extensive and I refer the interested reader to the literature [19]. The Yukawa couplings  $Qd^c h$ ,  $Qu^c \bar{h}$  and  $Le^c h$  are invariant under the custodial  $SU(2)_c$  symmetry. Therefore, the same fermion mass textures are expected to arise as in the models in which the custodial  $SU(2)$  is absent [6,20].

A similar mechanism may be possible in the case of superstring flipped  $SU(5)$  models [11,15] and other string GUT models [17]. For example, additional gauge symmetries from twisted sectors were shown to arise in the flipped  $SU(5)$  model

of Ref. [15]. The extended symmetries arise because of the existence of a sector in the additive group, of the form of Eq. (3), with  $X_L \cdot X_L = 0$ . Such a sector exist in the additive group because of the assignment of boundary conditions to the set of internal world-sheet fermions  $\{y, \omega|\bar{y}, \bar{\omega}\}^{1,\dots,6}$ . In terms of flipped  $SU(5)$  representations the dimension four operators arise from the operator

$$FF\bar{f}H \tag{11}$$

where  $F$  and  $H$  are in the 10 representation of  $SU(5)$  and  $\bar{f}$  is in the  $\bar{5}$  representation of  $SU(5)$ . The decomposition under standard model representations is:  $F = (Q, d_L^c, N_L^c)$ ,  $\bar{f} = (u_L^c, L)$ . The neutral state in  $H$  obtains a GUT scale VEV and breaks the  $SU(5) \times U(1)$  symmetry to  $SU(3)_C \times SU(2)_L \times U(1)_Y$ . However, such a VEV, in general, generates also the dangerous dimension four operators. In the presence of custodial nonabelian gauge symmetries, similar to the one shown to arise in some standard-like models, the dangerous operators may be forbidden due to the custodial symmetry.

Extended symmetries from twisted sectors may have additional phenomenological implications. In Ref. [21] extended gauge symmetries from twisted sectors were sought in type II superstring in order to circumvent the no go theorem of Ref. [22]. Examining the GSO phases in the superstring standard-like models it is observed that different choices of GSO phases result in different extensions of the gauge group. For example, in the model of Ref. [13] extended gauge symmetries may arise from the sector  $b_1 + b_2 + b_3 + \alpha + \beta + \gamma + (I)$ , where  $I = 1 + b_1 + b_2 + b_3$ . With the choice of GSO phases in Ref. [13] all the extra gauge bosons are projected out by the GSO projections. However, with the modified GSO phases

$$c \begin{pmatrix} 1 \\ \gamma \end{pmatrix} \rightarrow -c \begin{pmatrix} 1 \\ \gamma \end{pmatrix}, c \begin{pmatrix} \alpha \\ \beta \end{pmatrix} \rightarrow -c \begin{pmatrix} \alpha \\ \beta \end{pmatrix} \text{ and } c \begin{pmatrix} \gamma \\ \beta \end{pmatrix} \rightarrow -c \begin{pmatrix} \gamma \\ \beta \end{pmatrix},$$

additional space-time vector bosons are obtained from the sector  $b_1 + b_2 + b_3 + \alpha + \beta + \gamma + (I)$ . The sector  $b_1 + b_2 + b_3 + \alpha + \beta + \gamma + (I)$  produces the representations  $3_1 + 3_{-1}$

of  $SU(3)_H$ , where one of the  $U(1)$  combinations is the  $U(1)$  in the decomposition of  $SU(4)$  under  $SU(3) \times U(1)$ . In this case the hidden  $SU(3)_H$  gauge group is extended to  $SU(4)_H$ . Thus, the hidden sector contains two nonabelian factors  $SU(5) \times SU(4)$ . The possibility of extending the hidden sector gauge group from twisted sectors may be instrumental in trying to implement the dilaton stabilization mechanism of Ref. [23]. It should be noted that in the model of Ref. [14] the additional space-time vector bosons from the sector  $1 + \alpha + 2\gamma$  can be projected from the massless spectrum as well. However, if we require  $N = 1$  space-time supersymmetry then the projection of the additional gauge bosons is correlated with the presence of Higgs doublets from the sector  $b_1 + b_2 + \alpha + \beta$  in the massless spectrum. This is easily seen by substituting the vectors  $\alpha$  and  $\beta$  in the basis with the vectors  $b_1 + b_2 + \alpha + \beta$  and  $1 + \alpha + 2\gamma$ . It is observed that the intersection between these two vectors is empty. Therefore, the GSO projections correlate between Higgs doublets with extra gauge bosons or Higgs triplets without extra gauge bosons. The only other alternative that was found to project out the extra gauge bosons is by projecting out the last surviving gravitino, thus breaking supersymmetry at the Planck scale. This is achieved by modifying the phase  $c \begin{pmatrix} S \\ \alpha \end{pmatrix} \rightarrow -c \begin{pmatrix} S \\ \alpha \end{pmatrix}$ . It would be of interest to examine whether this class of nonsupersymmetric models can produce realistic phenomenology and vanishing cosmological constant. Finally, the extended gauge symmetries provide some freedom in the definition of the weak hypercharge. We may still define the weak hypercharge to be  $U(1)_Y = 1/3U(1)_C + 1/2U(1)_L$  with the standard  $SO(10)$  embedding. However, we may also define it as  $U(1)_Y = U(1)_{C'} + 1/2U(1)_L$ . With these two definitions, the weak hypercharge of the quark and leptons are the same. This freedom may be instrumental in trying to understand the disparity between the gauge coupling unification scale in the Minimal Supersymmetric Standard Model (MSSM) and the string unification scale.

In this paper I have shown how additional gauge bosons may appear in realistic superstring derived models from twisted sectors. In the free fermionic models

that are based on  $Z_2 \times Z_2$  orbifold with standard embedding the additional gauge bosons may give rise to custodial nonabelian gauge symmetries, under which only the leptons transform. In the fermionic models the extended symmetries arise because of the asymmetry of the boundary conditions in the vectors  $\alpha, \beta, \gamma$  between the left,  $\{y, \omega\}$ , and right  $\{\bar{y}, \bar{\omega}\}$ , internal world-sheet fermions. In the orbifold formulation the extended symmetries should be regarded as arising due to the asymmetry of the twists  $\alpha, \beta, \gamma$ . As a result of the custodial symmetry, dimension four baryon number violating operators are forbidden to all orders of nonrenormalizable terms while dimension four lepton number violating operators are allowed. Consequently, R-parity may be broken close to the Planck scale, but proton decay cannot be mediated by dimension four operators. It should be noted that the custodial symmetry also forbid the dimension five operator  $QQQL$ . Combined with the selection imposed by the left-moving global  $U(1)$  symmetries, this implies that dimension five operators that may result in proton decay are forbidden to all orders of nonrenormalizable terms. The effective low energy superpotential may contain the dimension four lepton number violating operator,  $QLd^c$ , while the dimension four baryon violating operator is forbidden. Due to the absence of continuous global symmetries in superstring theory [24] and possibly in any Planck scale theory, a mechanism similar in nature to the one proposed in this paper, may be the only possible avenue to avoid proton decay in supersymmetric Planck scale theories. The possible R parity violating terms are specified explicitly in specific models. Consequently, the specific low energy predictions are expected to be different from the low energy phenomenology of the MSSM.

## REFERENCES

1. For reviews see, H.P. Nilles, Phys. Rep. **110** (1984) 1; D.V. Nanopoulos and Lahanas, Phys. Rep. **145** (1987) 1.
2. M. Green, J. Schwarz and E. Witten, Superstring Theory, 2 vols., Cambridge University Press, 1987.
3. M.Dine, V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg, Nucl. Phys. B **259** (1985) 549.
4. B. Greene *et al.*, Phys.Lett.**B180** (1986) 69; Nucl.Phys.**B278** (1986) 667; **B292** (1987) 606; R. Arnowitt and P. Nath, Phys.Rev.**D39** (1989) 2006; **D42** (1990) 2498; Phys.Rev.Lett. **62** (1989) 222.
5. J. Ellis, J. Lopez and D.V. Nanopoulos, Phys. Lett. B **252** (1990) 53, G. Leontaris and T. Tamvakis, Phys. Lett. B **260** (1991) 333.
6. A.E. Faraggi, Nucl. Phys. B **403** (1993) 101, hep-th/9208023; IASSNS-94-18, hep-ph/9403312, Nucl. Phys. **B**, in press.
7. A. Font, L.E. Ibañez and F. Quevedo, Phys.Lett.**B228** (1989) 79.
8. A.E. Faraggi, D.V. Nanopoulos and K. Yuan, Nucl. Phys. B **335** (1990) 347.
9. A.E. Faraggi, Phys. Lett. B **245** (1990) 435; A.E. Faraggi and E. Halyo, Phys. Lett. B **307** (1993) 311, hep-th/9303060.
10. I. Antoniadis, C. Bachas, and C. Kounnas, Nucl.Phys.**B289** (1987) 87; I. Antoniadis and C. Bachas, Nucl.Phys.**B298** (1988) 586; H. Kawai, D.C. Lewellen, and S.H.-H. Tye, Nucl.Phys.**B288** (1987) 1; R. Bluhm, L. Dolan, and P. Goddard, Nucl.Phys.**B309** (1988) 330.
11. I. Antoniadis, J. Ellis, J. Hagelin, and D.V. Nanopoulos, Phys. Lett. B **231** (1989) 65; I. Antoniadis, G. K. Leontaris and J. Rizos, Phys. Lett. B **245** (1990) 161.
12. A.E. Faraggi, Nucl. Phys. B **387** (1992) 239, hep-th/9208024.

13. A.E. Faraggi, Phys. Lett. B **278** (1992) 131.
14. A.E. Faraggi, Phys. Lett. B **274** (1992) 47.
15. I. Antoniadis, J. Ellis, S. Kelley, and D.V. Nanopoulos, Phys. Lett. B **272** (1991) 31; J. Lopez, D.V. Nanopoulos and K. Yuan, Nucl. Phys. B **399** (1993) 654, hep-th/9203025.
16. A.E. Faraggi, Phys. Lett. B **326** (1994) 62, hep-ph/9311312.
17. S. Chaudhuri, S. Chung, and J. D. Lykken, FERMILAB-PUB-94-137-T,  
Talk given at 2nd IFT Workshop on Yukawa Couplings and the Origins of  
Mass, Gainesville, Feb 1994, hep-ph/9405374.
18. M. Dine, N. Seiberg and E. Witten, Nucl. Phys. B **289** (1987) 585.
19. There are numerous papers on this subject. A partial list includes L.J. Hall  
and M. Suzuki, Nucl. Phys. B **231** (1984) 419; I.H. Lee, Nucl. Phys. B **246**  
(1984) 120; S. Dawson, Nucl. Phys. B **261** (1985) 297 297; R. Barbieri and  
A. Masiero Nucl. Phys. B **267** (1986) 679; V. Barger, G.F. Giudice, and T.  
Han, Phys. Rev. D **40** (1989) 2987; S. Dimopoulos *et al.*, Phys. Rev. D **40**  
(1989) 2987; H. Dreiner and G.G. Ross, Nucl. Phys. B **365** (1991) 597.
20. A.E. Faraggi, Nucl. Phys. B **407** (1993) 57, hep-ph/9210256; A.E.  
Faraggi and E.Halyo, Phys. Lett. B **307** (1993) 305, hep-ph/9301261;  
Nucl. Phys. B **416** (1994) 63, hep-ph/9306235.
21. L. Dolan and S. Horvath , Nucl. Phys. B **416** (1994) 87.
22. L. Dixon, V. Kaplunovsky and C. vafa, Nucl. Phys. B **294** (1987) 43.
23. N.V. Krasnikov, Phys. Lett. B **193** (1987) 37; L. Dixon, in Proc. A.P.S. DPF  
Meeting, Houston, TX. 1990; T. Taylor, Phys. Lett. B **252** (1990) 59.
24. T. Banks and L. Dixon, Nucl. Phys. B **307** (1988) 93.

	$\psi^\mu$	$\{\chi^{12}; \chi^{34}; \chi^{56}\}$	$\bar{\psi}^1, \bar{\psi}^2, \bar{\psi}^3, \bar{\psi}^4, \bar{\psi}^5, \bar{\eta}^1, \bar{\eta}^2, \bar{\eta}^3$	$\bar{\phi}^1, \bar{\phi}^2, \bar{\phi}^3, \bar{\phi}^4, \bar{\phi}^5, \bar{\phi}^6, \bar{\phi}^7, \bar{\phi}^8$
$\alpha$	0	{0, 0, 0}	1, 1, 1, 0, 0, 0, 0, 0	1, 1, 1, 1, 0, 0, 0, 0
$\beta$	0	{0, 0, 0}	1, 1, 1, 0, 0, 0, 0, 0	1, 1, 1, 1, 0, 0, 0, 0
$\gamma$	0	{0, 0, 0}	$\frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}$	$\frac{1}{2}, 0, 1, 1, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, 0$

	$y^3 y^6, y^4 \bar{y}^4, y^5 \bar{y}^5, \bar{y}^3 \bar{y}^6$	$y^1 \omega^6, y^2 \bar{y}^2, \omega^5 \bar{\omega}^5, \bar{y}^1 \bar{\omega}^6$	$\omega^1 \omega^3, \omega^2 \bar{\omega}^2, \omega^4 \bar{\omega}^4, \bar{\omega}^1 \bar{\omega}^3$
$\alpha$	1, 1, 1, 0	1, 1, 1, 0	1, 1, 1, 0
$\beta$	0, 1, 0, 1	0, 1, 0, 1	1, 0, 0, 0
$\gamma$	0, 0, 1, 1	1, 0, 0, 0	0, 1, 0, 1

*Table 1.* A three generations standard-like model that produces a custodial  $SU(2)$  symmetry.

F	SEC	$SU(3)_C \times SU(2)_L \times SU(2)_c$	$Q_{C'}$	$Q_L$	$Q_1$	$Q_2$	$Q_3$	$Q_{4'}$	$Q_{5'}$	$SU(5) \times SU(3)$	$Q_{7''}$	$Q_8$
$L_1$	$b_1 \oplus$	(1,2,2)	$-\frac{1}{2}$	0	$\frac{1}{2}$	0	0	$-\frac{1}{2}$	$-\frac{1}{2}$	(1,1)	$-\frac{2}{3}$	0
$Q_1$	$1 + \alpha + 2\gamma$	(3,2,1)	$\frac{1}{6}$	0	$\frac{1}{2}$	0	0	$-\frac{1}{2}$	$-\frac{1}{2}$	(1,1)	$\frac{4}{3}$	0
$d_1$		( $\bar{3}$ ,1,1)	$-\frac{1}{6}$	1	$\frac{1}{2}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$-\frac{4}{3}$	0
$N_1$		(1,1,2)	$\frac{1}{2}$	-1	$\frac{1}{2}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$\frac{2}{3}$	0
$e_1$		(1,1,2)	$\frac{1}{2}$	1	$\frac{1}{2}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$\frac{2}{3}$	0
$u_1$		( $\bar{3}$ ,1,1)	$-\frac{1}{6}$	-1	$\frac{1}{2}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$-\frac{4}{3}$	0
$L_2$	$b_2 \oplus$	(1,2,2)	$-\frac{1}{2}$	0	0	$\frac{1}{2}$	0	$\frac{1}{2}$	$-\frac{1}{2}$	(1,1)	$-\frac{2}{3}$	0
$Q_2$	$1 + \alpha + 2\gamma$	(3,2,1)	$\frac{1}{6}$	0	0	$\frac{1}{2}$	0	$\frac{1}{2}$	$-\frac{1}{2}$	(1,1)	$\frac{4}{3}$	0
$d_2$		( $\bar{3}$ ,1,1)	$-\frac{1}{6}$	1	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$-\frac{4}{3}$	0
$N_2$		(1,1,2)	$\frac{1}{2}$	-1	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$\frac{2}{3}$	0
$e_2$		(1,1,2)	$\frac{1}{2}$	1	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$\frac{2}{3}$	0
$u_2$		( $\bar{3}$ ,1,1)	$-\frac{1}{6}$	-1	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$-\frac{4}{3}$	0
$L_3$	$b_3 \oplus$	(1,2,2)	$-\frac{1}{2}$	0	0	0	$\frac{1}{2}$	0	1	(1,1)	$-\frac{2}{3}$	0
$Q_3$	$1 + \alpha + 2\gamma$	(3,2,1)	$\frac{1}{6}$	0	0	0	$\frac{1}{2}$	0	1	(1,1)	$\frac{4}{3}$	0
$d_3$		( $\bar{3}$ ,1,1)	$-\frac{1}{6}$	1	0	0	$\frac{1}{2}$	0	-1	(1,1)	$-\frac{4}{3}$	0
$N_3$		(1,1,2)	$\frac{1}{2}$	-1	0	0	$\frac{1}{2}$	0	-1	(1,1)	$\frac{2}{3}$	0
$e_3$		(1,1,2)	$\frac{1}{2}$	1	0	0	$\frac{1}{2}$	0	-1	(1,1)	$\frac{2}{3}$	0
$u_3$		( $\bar{3}$ ,1,1)	$-\frac{1}{6}$	-1	0	0	$\frac{1}{2}$	0	-1	(1,1)	$-\frac{4}{3}$	0

Table 2. Massless states and their quantum numbers in the model of table 1.



F	SEC	$SU(3)_C \times SU(2)_L \times SU(2)_c$	$Q_{C'}$	$Q_L$	$Q_1$	$Q_2$	$Q_3$	$Q_{4'}$	$Q_{5'}$	$SU(5) \times SU(3)$	$Q_{7''}$	$Q_8$
$V_1$	$b_1 + 2\gamma$	(1,1,1)	$\frac{1}{4}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	(1,3)	$-\frac{4}{3}$	$\frac{5}{2}$
$\bar{V}_1$		(1,1,1)	$-\frac{1}{4}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	(1, $\bar{3}$ )	$\frac{4}{3}$	$-\frac{5}{2}$
$T_1$		(1,1,1)	$\frac{1}{4}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	(5,1)	$-\frac{4}{3}$	$-\frac{3}{2}$
$\bar{T}_1$		(1,1,1)	$-\frac{1}{4}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	( $\bar{5}$ ,1)	$\frac{4}{3}$	$\frac{3}{2}$
$V_2$	$b_2 + 2\gamma$	(1,1,1)	$\frac{1}{4}$	0	$\frac{1}{2}$	0	$\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	(1,3)	$-\frac{4}{3}$	$\frac{5}{2}$
$\bar{V}_2$		(1,1,1)	$-\frac{1}{4}$	0	$\frac{1}{2}$	0	$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	(1, $\bar{3}$ )	$\frac{4}{3}$	$-\frac{5}{2}$
$T_2$		(1,1,1)	$\frac{1}{4}$	0	$\frac{1}{2}$	0	$\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	(5,1)	$-\frac{4}{3}$	$-\frac{3}{2}$
$\bar{T}_2$		(1,1,1)	$-\frac{1}{4}$	0	$\frac{1}{2}$	0	$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	( $\bar{5}$ ,1)	$\frac{4}{3}$	$\frac{3}{2}$
$V_3$	$b_3 + 2\gamma$	(1,1,1)	$\frac{1}{4}$	0	$\frac{1}{2}$	$\frac{1}{2}$	0	0	-1	(1,3)	$-\frac{4}{3}$	$\frac{5}{2}$
$\bar{V}_3$		(1,1,1)	$-\frac{1}{4}$	0	$\frac{1}{2}$	$\frac{1}{2}$	0	0	1	(1, $\bar{3}$ )	$\frac{4}{3}$	$-\frac{5}{2}$
$T_3$		(1,1,1)	$\frac{1}{4}$	0	$\frac{1}{2}$	$\frac{1}{2}$	0	0	-1	(5,1)	$-\frac{4}{3}$	$-\frac{3}{2}$
$\bar{T}_3$		(1,1,1)	$-\frac{1}{4}$	0	$\frac{1}{2}$	$\frac{1}{2}$	0	0	1	( $\bar{5}$ ,1)	$\frac{4}{3}$	$\frac{3}{2}$
$D_1$	$b_1 + b_3 + \beta \pm \gamma$	(3,1,1)	$\frac{5}{24}$	$\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	0	0	(1,1)	0	$-\frac{15}{4}$
$\bar{D}_1$		( $\bar{3}$ ,1,1)	$-\frac{5}{24}$	$-\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{4}$	0	0	(1,1)	0	$\frac{15}{4}$
$H_1$		(1,1,1)	$-\frac{5}{8}$	$\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	0	0	(1, $\bar{3}$ )	0	$\frac{5}{4}$
$\bar{H}_1$		(1,1,1)	$\frac{5}{8}$	$-\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{4}$	0	0	(1,3)	0	$-\frac{5}{4}$
$D_2$	$b_2 + b_3 + \beta \pm \gamma$	(3,1,1)	$\frac{5}{24}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	0	0	(1,1)	0	$-\frac{15}{4}$
$\bar{D}_2$		( $\bar{3}$ ,1,1)	$-\frac{5}{24}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	0	0	(1,1)	0	$\frac{15}{4}$
$H_2$		(1,1,1)	$-\frac{5}{8}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	0	0	(1, $\bar{3}$ )	0	$\frac{5}{4}$
$\bar{H}_2$		(1,1,1)	$\frac{5}{8}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	0	0	(1,3)	0	$-\frac{5}{4}$
$\ell_1$	$1 + b_1 + \alpha + 2\gamma$	(1,2,1)	$\frac{1}{2}$	0	$-\frac{1}{2}$	0	0	$-\frac{1}{2}$	$-\frac{1}{2}$	(1,1)	$-\frac{8}{3}$	0
$S_1$		(1,1,1)	$-\frac{1}{2}$	-1	$-\frac{1}{2}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$\frac{8}{3}$	0
$S'_1$		(1,1,1)	$-\frac{1}{2}$	1	$-\frac{1}{2}$	0	0	$\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$\frac{8}{3}$	0
$\ell_2$	$1 + b_2 + \alpha + 2\gamma$	(1,2,1)	$\frac{1}{2}$	0	0	$-\frac{1}{2}$	0	$\frac{1}{2}$	$-\frac{1}{2}$	(1,1)	$-\frac{8}{3}$	0
$S_2$		(1,1,1)	$-\frac{1}{2}$	-1	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$\frac{8}{3}$	0
$S'_2$		(1,1,1)	$-\frac{1}{2}$	1	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	$\frac{1}{2}$	(1,1)	$\frac{8}{3}$	0
$\ell_3$	$1 + b_3 + \alpha + 2\gamma$	(1,2,1)	$\frac{1}{2}$	0	0	0	$-\frac{1}{2}$	0	1	(1,1)	$-\frac{8}{3}$	0
$S_3$		(1,1,1)	$-\frac{1}{2}$	-1	0	0	$-\frac{1}{2}$	0	-1	(1,1)	$\frac{8}{3}$	0
$S'_3$		(1,1,1)	$-\frac{1}{2}$	1	0	0	$-\frac{1}{2}$	0	-1	(1,1)	$\frac{8}{3}$	0

*Table 3.* Massless states and their quantum numbers in the model of table 1.